

Inflation in large N limit of supersymmetric gauge theories

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Abstract

Within supersymmetry we provide an example where the inflaton sector is derived from a gauge invariant polynomial of $SU(N)$ or $SO(N)$ gauge theory. Inflation in our model is driven by multi-flat directions, which assist accelerated expansion. We show that multi-flat directions can flatten the individual non-renormalizable potentials such that inflation can occur at sub-Planckian scales. We calculate the density perturbations and the spectral index, we find that the spectral index is closer to scale invariance for large N . In order to realize a successful cosmology we require large N of order, $N \sim 600$.

1 Introduction

The current satellite based experiments suggest that the early universe might have had a spell of accelerated expansion [1]. The idea of inflation is interesting as it solves quite a range of problems [2]. Usually in the literature it is assumed that a single gauge singlet with its non-vanishing potential energy is responsible for driving inflation, see [3]. However a gauge singlet is a step backward in presenting cosmological inflationary models, because the gauge couplings and the masses can be tweaked at our will, or from the observational constraints, and not from any fundamental gauged sector ¹. The lack of motivation behind the absolute gauge singlet in nature enforces gauging the inflaton sector. It is equally appealing if we wish to connect the inflaton sector to a realistic particle physics model.

Supersymmetric gauge theories bring a new host of flat directions along which the scalar potential vanishes. Within $N = 1$ SUSY it implies that D-term and F-term vanishes identically. These flat directions are derived from gauge invariant polynomials

¹In some models of inflation the Higgs potential is responsible for providing the non-zero vacuum expectation value during inflation. However inflation in these models can last forever if there is no first or second order phase transition. The phase transition leads to an end of inflation due to a rolling scalar field which is usually treated as a gauge singlet inflaton [4]. This particular feature is explicit in a supersymmetric set up where the total potential, within $N = 1$ supersymmetry (SUSY), is obtained from the superpotential, $W = \lambda\Phi(H^2 - H_0)^2$, where Φ is the absolute gauge singlet and H is the Higgs sector.

[5–7], which are build on gauge invariant combination of squarks and sleptons. Interestingly SUSY flat directions do not receive any perturbative corrections [8]. However they can obtain non-perturbative superpotential and Kähler corrections when supersymmetry is broken [6, 9]. These flat directions have many cosmological implications, such as baryogenesis [6], dark matter [10], and as a potential source for the cosmic density perturbations [11], for a review on the subject, see [12]. In this paper our main goal is to seek a gauge invariant flat direction(s) as a candidate for the inflaton within minimal supersymmetric standard model (MSSM) and/or beyond MSSM.

We will explain the caveats in treating MSSM flat directions as an inflaton. MSSM flat directions are usually parameterized by a monomial, which does not lead to any inflation because non-renormalizable potentials dominate at large vacuum expectation values (vevs), whose contributions can only be trusted below their cut-off. It is helpful to represent the flat direction by gauge invariant polynomials. This brings a new host of flat directions, see [13]. We will argue that spanning the moduli space of flat directions can lead to inflation due to a collective motion of these flat directions in a moduli space.

The concept of having inflation from multi-fields is well known. In Ref. [14], it was first demonstrated that inflation occurs at ease with many exponentials, and it has been generalized to many other forms of potentials, see [15].

In this paper we will argue that enhancing the gauge group along with the matter content will provide multi-flat directions within $SU(N)$ and/or $SO(N)$ gauge theories, where inflation will be driven collectively by the degrees of freedom below the cut-off scale, which we consider here as the Planck scale.

2 Why is it hard to obtain inflation from MSSM flat directions?

Within MSSM there are many flat directions subject to F and D constraints

$$F_i \equiv \frac{\partial W}{\partial \Phi_i} = 0, \quad D^A \equiv \Phi^\dagger T^A \Phi = 0, \quad (1)$$

for the scalars Φ_i . More elegant way of describing a flat direction is through gauge invariant holomorphic polynomials of the chiral superfields Φ_i . Within MSSM, with R-parity, all the flat directions have been tabulated, see [7]. In a cosmological context where supersymmetry is broken by the finite energy density of the Universe, the flat directions obtain various corrections. Notably the non-renormalizable superpotential corrections are of types [6, 8]

$$W \sim \lambda \frac{\Phi^n}{M_p^{n-3}}, \quad W \sim \lambda_1 \frac{\Phi^{n-1} \Psi}{M_p^{n-3}}, \quad (2)$$

where, Φ, Ψ are flat directions, $\lambda, \lambda_1 \sim \mathcal{O}(1)$, and $M_p \sim 10^{18}$ GeV. Within MSSM, with R-parity, all the flat directions are lifted by the non-renormalizable operators, $n = 4, 5, 6, 7, 9$, see [7].

Besides such a non-renormalizable correction, the MSSM flat directions naturally obtain soft SUSY breaking mass terms,

$$V \sim m_{soft}^2 \Phi^2, \quad (3)$$

where $m_{soft} \sim \mathcal{O}(1)$ TeV. Note that the above potential is similar to that of a chaotic type potential [16], but with a mass parameter which is way too small to provide any observable effects, since $\delta\rho/\rho \sim m_{soft}/M_p \sim \mathcal{O}(\text{TeV}/M_p) \sim 10^{-15}$. Besides COBE/WMAP normalization requires that $m_{soft} \sim 10^{13}$ GeV [17], which is ten orders of magnitude larger than what we expect from soft SUSY breaking mass term.

At sufficiently large vevs, $\Phi, \Psi \gg \mathcal{O}(\text{TeV})$, the potential from the non-renormalizable superpotential terms dominate over the soft SUSY breaking potential. One might expect to realize inflation from the MSSM flat directions with a potential derived from Eq. (3),

$$V \sim |\lambda|^2 \frac{\Phi^{2(n-1)}}{M_p^{2n-6}}. \quad (4)$$

The above potential mimics that of a chaotic type inflationary model [16]. Inflation with such a potential is possible only when the vev is larger than the cut-off. Inflation ends when $\Phi_{end} \sim (2n-2)M_p$. However in our case, since Φ is a gauge invariant quantity, we can no longer trust the non-renormalizable potential above the cut-off with a coefficient, $\lambda \sim \mathcal{O}(1)$. Further note that even if we take the vev larger than M_p , the potential is too steep to give rise an interesting effect. Furthermore the coupling constant should be $\lambda \lesssim 10^{-14}$ [17], in order to provide the right amplitude for the density perturbations, which is ridiculously small compared to any SM/MSSM Yukawa and/or gauge couplings. These facts immediately suggest that inflation with the MSSM flat directions is very unlikely. The main challenge is to realize inflation at vevs smaller than the Planck scale, which is impossible to realize with a single flat direction.

Besides the superpotential contribution, the flat direction also obtains correction due to the Kähler potential in $N=1$ supergravity. For a minimal choice of Kähler potential, $K = \pm\Phi_i^\dagger\Phi_i + \dots$, the flat direction obtains a potential,

$$V \sim \pm\mathcal{O}(1)H^2\Phi^2, \quad (5)$$

where H is the Hubble rate during and after inflation. However there is a bit of freedom in the choice of a Kähler potential. Usually such corrections are not present in No-Scale type kähler models [18]. For our purpose we will ignore such a dangerous correction which leads to an unsuccessful inflationary scenario. This is because a positive Hubble induced mass correction will spoil any interesting dynamics.

Motivated by these problems we are inclined to study the properties of multi flat directions, which will evolve collectively in such a way to assist inflation similar to the assisted inflationary scenarios [14]². One of the revealing properties of the assisted

²If many scalar fields evolve independently then they assist inflation inspite of the fact that individual potential is unable to sustain inflation on its own. The key point is that collective dynamics of fields increase the Hubble friction term which leads to the slow rolling of the fields. The main constraint is that the fields ought to be free, the coupling between the fields do not assist inflation [14, 15].

inflation is that it is possible to drive inflation for chaotic type potentials, $V \sim \sum_i \phi_i^n$, with vevs below the Planckian scale, see also [19]. This is a good news for us, because the non-renormalizable potentials are trusted only below the cut-off. Therefore there is a glimmer of hope that we might be able to sustain inflation from the flat directions with potentials of type Eq. (4), albeit with many directions.

However the question is to seek whether MSSM can provide us with sufficient number of independent flat directions? For example, LH_u , udd , directions are lifted by $n = 4$ non-renormalizable operators, which can obtain large vevs simultaneously, but LH_u , LLe directions, depending on the family indices, need not be simultaneously flat. As we shall show in section IV we would require at least 600 independent directions. Unfortunately within MSSM there are only 334 D-flat directions. This number is further reduced due to the F-term constraints, and moreover these directions are not independent to each other.

So much for a flat direction represented by monomials. Very recently we explored, for the first time, the dynamics of multiple flat directions parameterized by the gauge invariant polynomials in Ref. [13]. It turns out that the multiple flat directions have very interesting properties. We mention them briefly; for a single flat direction the choice of a gauge is trivially satisfied to be a pure gauge by definition, see for detailed discussion in Ref. [13]. However the choice of a pure gauge is not trivially satisfied in the multi-flat direction case, unless one assumes that the gauge degrees of freedom are frozen with a homogeneous distribution without any spatial perturbations. Any spatial fluctuation will excite the gauge degrees of freedom with an interesting astrophysical implication, such as exciting the seed magnetic field, see [20].

In Ref. [13], we studied a polynomial I spanned by the Higgses and the sleptons,

$$I = \nu_1 H_u L_1 + \nu_2 H_u L_2 + \nu_3 H_u L_3 \quad (6)$$

where ν_i are complex coefficients, H_u is the up-type Higgs and L_i are the sleptons. For a following field configuration, the polynomial I has a vanishing matter current and vanishing gauge fields [13],

$$L_i = e^{-i\chi/2} \phi_i \begin{pmatrix} 0 \\ 1 \end{pmatrix}, \quad H_u = e^{i\chi/2} \sqrt{\sum_i |\phi_i|^2} \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad (7)$$

where ϕ_i are complex scalar fields and the phase χ is a real field constrained by

$$\partial_\mu \chi = \frac{\sum_j J_j^\phi}{2i \sum_k |\phi_k|^2}, \quad J_i^\phi = \phi_i^* \partial_\mu \phi_i - \phi_i \partial_\mu \phi_i^*. \quad (8)$$

The field configuration in Eq. (7) leads to an effective Lagrangian for the flat direction fields ϕ_i ,

$$\mathcal{L} = |D_\mu H_u|^2 + \sum_{i=1}^3 |D_\mu L_i|^2 - V = \frac{1}{2} \partial_\mu \Phi^\dagger \left(1 + P_1 - \frac{1}{2} P_2 \right) \partial^\mu \Phi - V, \quad (9)$$

where D_μ is a gauge covariant derivative that reduces to the partial derivative when the gauge fields vanish, P_1 is the projection operator along Φ and P_2 along Ψ , where

$$\bar{\phi} = (\phi_1 \quad \phi_2 \quad \phi_3)^T, \quad \Phi = \begin{pmatrix} \bar{\phi} \\ \phi^* \end{pmatrix}, \quad \Psi = \begin{pmatrix} \bar{\phi} \\ -\bar{\phi}^* \end{pmatrix}, \quad (10)$$

and the corresponding equation of motion

$$\begin{aligned} \partial_\mu \partial^\mu \Phi + 3H\dot{\Phi} + \left(1 - \frac{1}{2}P_1 + P_2\right) \frac{\partial V}{\partial \Phi^\dagger} - R^{-2} \left[\partial_\mu \Psi (\Psi^\dagger \partial^\mu \Phi) \right. \\ \left. + \Psi (\partial_\mu \Psi^\dagger P_2 \partial^\mu \Phi) + \frac{1}{2} \Phi \partial_\mu \Phi^\dagger \left(1 - P_1 - \frac{3}{2}P_2\right) \partial^\mu \Phi \right] = 0, \end{aligned} \quad (11)$$

where $R = \sqrt{\Phi^\dagger \Phi}$. We are interested in the background dynamics where all the fields are homogeneous in time, and for simplicity we study only the radial motion, such that $\Phi = R\hat{e}_\Phi$, where $\dot{\hat{e}}_\Phi = 0$ (the dot denotes derivative w.r.t time). Then the equation of motion simplifies to

$$\ddot{R} + 3H\dot{R} + \frac{1}{2} \frac{\partial V}{\partial R} = 0, \quad (12)$$

if we assume $V = V(R)$ for simplicity. A notable feature is that the fields have non-minimal kinetic terms, since the field manifold defined by the flat direction is curved, actually a hyperbolic manifold. This results into the usual equation of motion for one scalar field with a potential for the radial mode except for the factor 1/2, which makes the potential effectively flatter in this direction. This can be traced to the square root nature of H_u in Eq. (7).

This illustrates that there is a way of making the flat direction potential even flatter in the equation of motion. Now one might, of course, wonder whether MSSM flat directions might work. As far as our example of LH_u is concerned there are only three families which we can account for. The flattest MSSM direction, $QuQue$, is lifted by $n = 9$ superpotential operator, $QuQuQuH_d ee$. The flat direction $QuQue$ is an 18 complex dimensional manifold as can be seen from Table 5 of Ref. [7]. The largest D-flat direction is only 33 complex dimensional [7]. As we shall see in section IV, we would require much larger number of fields to assist inflation below the Planck vev. The conclusion is that we must go beyond the MSSM gauge group to seek a large number of flat directions.

3 Construction of flat directions in $SU(N)$ and $SO(N)$ gauge theories

Now we wish to study the large N gauge group which could provide us with multi flat directions. We do not pretend here to scan all the D- and F- flat directions of either $SU(N)$ or $SO(N)$, but we take a particular combination which is indeed a flat direction. Encouraged by our previous study within MSSM, we consider M fields H_i and $N-1$ fields G_j in the fundamental representation N of the gauge group. Note that the matter content is also enhanced, which has a total $N-1+M$ degrees of freedom. Then there exists a D-flat direction described by a gauge invariant polynomial

$$I = \sum_{j=1}^M \alpha_j \epsilon_{d_1 \dots d_{N-1}} H_1^{d_1} \dots H_{N-1}^{d_{N-1}} G_j^e, \quad (13)$$

which after solving the constraint equations

$$\frac{\partial I}{\partial H_j^a} = CH_j^{a*}, \quad \frac{\partial I}{\partial G_i^a} = CG_i^{a*} \quad (14)$$

produces a vacuum configuration

$$\begin{aligned} H_j^a &= \delta_N^a \phi_j & j = 1, \dots, M, \\ G_i^a &= \delta_i^a \sqrt{\sum_{j=1}^M |\phi_j|^2} & i = 1, \dots, N-1. \end{aligned} \quad (15)$$

When one substitutes Eq. (15) into D-terms one finds that all D-terms vanish.

In reality the solution of Eq. (14) is a gauge invariant surface in field space. In Eq. (15) we have chosen field values along this flat direction surface, which corresponds to fixing a particular gauge. Now we have to solve the corresponding gauge field value from the condition of vanishing matter current, and check whether it is indeed of pure gauge form. The matter current is given by

$$\begin{aligned} J_\mu^A &= ig \left[\sum_{j=1}^M \left(H_j^\dagger T^A D_\mu H_j - (D_\mu H_j)^\dagger T^A H_j \right) \right. \\ &\quad \left. + \sum_{i=1}^{N-1} \left(G_i^\dagger T^A D_\mu G_i - (D_\mu G_i)^\dagger T^A G_i \right) \right] \end{aligned} \quad (16)$$

where the covariant derivative is $D_\mu = \partial_\mu - igT^B A_\mu^B$, g is the gauge coupling constant, A_μ^B the gauge fields, and T^A are the gauge group generators in the defining representation, which are normalized as; $\text{Tr}(T^A T^B) = \frac{1}{2} \delta^{AB}$.

Substituting Eq. (15) into the matter current Eq. (16), we obtain

$$J_\mu^A = igT_{NN}^A \sum_{j=1}^M (\phi_j^* \partial_\mu \phi_j - \phi_j \partial_\mu \phi_j^*) + g^2 A_\mu^A \sum_{j=1}^M |\phi_j|^2 = 0. \quad (17)$$

Now it is a simple matter to solve the gauge field from this. In order to check whether it is a pure gauge, we must calculate the field strength tensor

$$F_{\mu\nu}^A = \partial_\mu A_\nu^A - \partial_\nu A_\mu^A + gf^{ABC} A_\mu^B A_\nu^C. \quad (18)$$

Since in the background the scalar fields are function of time alone, we have only $A_0(t) \neq 0$, rest of the gauge field components vanish. Therefore F_{00}^A is the only non-vanishing component, but this also vanishes due to anti-symmetry properties, so we have a pure gauge configuration, $F_{\mu\nu}^A = 0$, as required.

We also assume that the superpotential is such that Eq. (15) is also F-flat, although this is not really necessary, since any superpotential contributions can be taken into account in the potential.

4 Dynamics

For the purpose of illustration, we keep the potential for the flat direction, $V \sim f(|\Phi|^n/M_p^{n-4})$, though we will discuss the dynamics in a very general context. The Lagrangian for the flat direction is given by

$$\mathcal{L} = c \sum_{j=1}^M |D_\mu H_j|^2 + c \sum_{i=1}^{N-1} |D_\mu G_i|^2 - V(\{H_i, G_j\}), \quad (19)$$

where $c = 1/2$ for the real fields, and $c = 1$ for the complex fields. By inserting Eq. (15) into Eq. (19) leads to the Lagrangian³

$$\mathcal{L} = \frac{1}{2} \partial_\mu \Phi^\dagger [1 + (N-1)P] \partial^\mu \Phi - V(\Phi), \quad (20)$$

where $P = \Phi \Phi^\dagger / (\Phi^\dagger \Phi)$ is the projection operator. The field configurations of the real and the complex fields are

$$\Phi = (\phi_1, \dots, \phi_M)^T, \quad \phi_i \in \mathbb{R}, \quad (21)$$

$$\Phi = (\phi_1, \dots, \phi_M, \phi_1^*, \dots, \phi_M^*)^T, \quad \phi_i \in \mathbb{C}. \quad (22)$$

This Lagrangian results into an equation of motion

$$\begin{aligned} \partial_\mu \partial^\mu \Phi + 3H\dot{\Phi} + \frac{N-1}{N} \frac{\Phi}{\Phi^\dagger \Phi} |(1-P)\partial_\mu \Phi|^2 \\ + \left(1 - \frac{N-1}{N} P\right) \frac{\partial V}{\partial \Phi^\dagger} = 0, \end{aligned} \quad (23)$$

and the Friedmann equation

$$\begin{aligned} 3M_p^2 H^2 = \rho = \frac{1}{2} \dot{\Phi}^\dagger [1 + (N-1)P] \dot{\Phi} + V \\ + \frac{1}{2a^2} \partial_i \Phi^\dagger [1 + (N-1)P] \partial_i \Phi, \end{aligned} \quad (24)$$

where H is the Hubble parameter and a is the scale factor.

Now let us concentrate on the radial motion during inflation. Then it follows that the centrifugal acceleration vanishes in Eq. (23), and we can write down the equations of motion for $R = \sqrt{\Phi^\dagger \Phi}$ alone,

$$\ddot{R} + 3H\dot{R} + \frac{1}{N} V'(R) = 0, \quad (25)$$

$$3M_p^2 H^2 = \frac{N}{2} \dot{R}^2 + V(R). \quad (26)$$

Now the slow-roll approximation requires that; $|\ddot{R}| \ll 3H|\dot{R}|$, $|V'(R)/N|$ and $N\dot{R}^2/2 \ll V(R)$, so the equations simplify

$$\dot{R} \approx -\frac{V'(R)}{3NH}, \quad (27)$$

$$H \approx \sqrt{\frac{V(R)}{3M_p^2}}. \quad (28)$$

³Note that we have left out the gauge field contribution, since it only affects the dynamics transverse to the radial motion.

An equivalent dynamics can be obtained through the effective slow-roll parameters

$$\epsilon_{eff} = \frac{\epsilon}{N} \ll 1, \quad |\eta_{eff}| = \frac{|\eta|}{N} \ll 1, \quad (29)$$

where $\epsilon \equiv (M_p^2/2)(V'(R)/V(R))^2$ and $\eta \equiv M_p^2(V''(R)/V(R))$ are the usual slow roll parameters. Note an interesting point, both ϵ_{eff} and $|\eta_{eff}|$ become less than one for a large number of fields *i.e.* when $N \gg 1$. We can also define the number of e-foldings, \mathcal{N} , as

$$\mathcal{N} = \log \left(\frac{a(t_e)}{a(t)} \right) = \int_t^{t_e} H dt = \frac{N}{\sqrt{2}M_p} \int_{R_e}^R dR \frac{1}{\sqrt{\epsilon}}, \quad (30)$$

where the end of inflation is defined by, $\epsilon_{eff}(R_e) = 1$, and R_e is the value of R at the end of inflation. This shows that even if the potential were not flat enough for the usual slow-roll conditions to hold, the effective slow-roll parameters are small enough, therefore the required number of e-foldings can be generated if there are just enough many G type fields. In order to solve the flatness and the homogeneity problems we require $\mathcal{N} \sim 60$ e-foldings of inflation. This can be achieved by

$$N \sim 600 \left(\frac{\mathcal{N}}{60} \right) \left(\frac{\Delta R}{0.1M_p} \right)^{-1} \left(\frac{\epsilon}{2} \right)^{1/2}, \quad (31)$$

where $\Delta R \equiv R_i - R_e$. It is evident that we require large number of fields in order to realize inflation at vevs lower than M_p .

There are couple of points to be mentioned. First of all it is a good news that we can drive inflation with gauge invariant flat directions at vevs smaller than the cut-off scale, this is just one simple example we were looking for, however, the cost we have to pay is the large N group. Note here that N is not the dimension of the flat directions, but it is the number of colours in $SU(N)$, $SO(N)$ gauge theories. Further note that the matter content, $N - 1 + M$, in our case surpasses the number of colours N .

Our solution is still far from realistic. Nevertheless we take an important message from this analysis; perhaps it is extremely difficult to seek inflaton sector within a realistic, phenomenologically interesting, gauge group such as grand unified groups, e.g. $SU(5)$, $SO(10)$, etc.

5 Density Perturbations

For the sake of completeness we briefly discuss the density perturbation arising from the multi-fields. It is well known how to calculate the spectrum of the curvature perturbations in the case multiple fields, and for a non-flat field metric [21]

$$\mathcal{P}_{\mathcal{R}}(k) = \left(\frac{H}{2\pi} \right)^2 (G^{-1})_{ij} \frac{\partial \mathcal{N}}{\partial \phi_i} \frac{\partial \mathcal{N}}{\partial \phi_j}, \quad (32)$$

where G is the field metric and the summation goes over field components running through both ϕ_i and ϕ_i^* in the case of complex scalar fields. The isocurvature perturbations have been ignored in the above analysis. The inverse field metric is given

by

$$G^{-1} = 1 - \frac{N-1}{N} P. \quad (33)$$

Now assuming that the field trajectory is radial, the number of e-folds, \mathcal{N} , depends only on the radial motion. From Eqs. (30,32), the curvature perturbation spectrum is given by

$$\mathcal{P}_{\mathcal{R}}(k) = \frac{1}{N} \left(\frac{H}{\dot{R}} \right)^2 \left(\frac{H}{2\pi} \right)^2 = \frac{1}{24\pi^2 M_p^4 \epsilon_{eff}} \frac{V}{\epsilon_{eff}}. \quad (34)$$

Note that this last expression contains only the radial mode and it is the same as the usual formula for a single field case. This reiterates a point that the dynamics of multi-fields do not alter the spectrum of density perturbations, see [14].

Of particular interest is the spectral index n . This is given in our case by [21]

$$n - 1 = -6\epsilon_{eff} + 2\eta_{eff} = -\frac{6\epsilon - 2\eta}{N}. \quad (35)$$

This result shows that for larger N , the spectral index is even closer to the scale invariant.

During inflation there will be fluctuations along the transverse direction, with an amplitude $\sim H/2\pi$. There are $M - 1$ such modes along which the perturbations will give rise to isocurvature type fluctuations, whose analysis goes beyond the scope of the present paper.

6 Conclusion

In this letter we addressed the core issue of the inflaton sector; can we really find a workable model of inflation where the inflaton is *not* a gauge singlet? The answer is positive. Inflation can be driven by the gauge invariant multi-flat directions. Inflaton in our case is borne out of a gauge invariant quantity under $SU(N)$ and/or $SO(N)$ type SUSY gauge theories. We are also able to show that inflationary scale is sub-Planckian, however it requires a large N of order 600. This is not a very encouraging news, because as it appears from our analysis we require a large number of independent flat directions, which can be obtained only by increasing the number of colours, N , and the matter content.

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